

CERN-TH/96-237

USC-96/022

hep-th/9608183

Exceptional SW Geometry from ALE Fibrations

W. Lerche

Theory Division, CERN, 1211 Geneva 23, Switzerland

and

N.P. Warner

Physics Department, U.S.C., University Park, Los Angeles, CA 90089

We show that the genus 34 Seiberg-Witten curve underlying $N = 2$ Yang-Mills theory with gauge group E_6 yields physically equivalent results to the manifold obtained by fibration of the E_6 ALE singularity. This reconciles a puzzle raised by $N = 2$ string duality.

CERN-TH/96-237

August 1996

1. Introduction

There is now an extensive body of literature on the construction of the quantum effective actions on the Coulomb branch of $N = 2$ supersymmetric Yang-Mills theories [1]. The work of [1] was extended to other gauge groups: for example, to $G = SU(n)$ in [2,3], to $G = SO(2n+1)$ in [4], and to $G = SO(2n)$ in [5]. In these approaches, the curves are given in terms of the appropriate simple singularities W_{ADE} [6], and are generically of the form

$$y^2 = W^2(x; u_j) - \mu^2, \quad (1.1)$$

with $\mu = \Lambda^{h^\vee}$, where h^\vee is the dual Coxeter number of G , and Λ is the quantum scale. The $u_j, j = 1, \dots, \ell$ are the fundamental Casimir invariants (with degree increasing with the subscript j) and ℓ is the rank of G ; the top Casimir, u_ℓ , has degree h^\vee . For example, for $G = SU(n)$, one has $W_{A_{n-1}}(x; u_j) = x^n - \sum_{j=1}^{n-1} u_j x^{n-1-j}$.

In a complementary, unifying approach based on integrable systems [7], a general scheme for obtaining Seiberg-Witten (SW) curves for all groups was presented in [8]. As explained in more detail below, these curves are of the form

$$\zeta + \frac{\mu^2}{\zeta} + P_{\mathcal{R}}(x; u_j) = 0, \quad (1.2)$$

where $P_{\mathcal{R}}$ is a polynomial in x of order $\dim(\mathcal{R})$, where \mathcal{R} is some representation of G . For $G = SU(n)$, one can take $P_{\mathcal{R}}(x; u_j) \equiv W_{A_{n-1}}(x; u_j)$ so that the curves (1.2) and (1.1) are manifestly the same, up to a simple reparametrization. This however cannot be done for other groups, since $\dim(\mathcal{R})$ will in general¹ not match the degree of W_{ADE} , which is equal to h^\vee .

More recently, it was found in [9] how, via $N = 2$ heterotic-type II string duality [10], local SW geometry can be derived from fibrations of ALE spaces: the relevant manifolds are described by

$$\zeta + \frac{\mu^2}{\zeta} + W_{ADE}^{ALE}(x_1, x_2, x_3; u_j) = 0, \quad (1.3)$$

where W_{ADE}^{ALE} is the (non-compact) ALE space of type ADE ; for $G = SU(n)$, $W_{A_{n-1}}^{ALE} \equiv W_{A_{n-1}}(x_1; u_j) + x_2^2 + x_3^2$. Obviously, by trivially integrating out the quadratic pieces in x_2 and x_3 (which does not change the singularity type) this manifold is equivalent to the above SW curves. It also gives rise to the same periods [9].

¹ However, this can easily be reconciled [5] for $G = SO(2n)$.

Actually, by simple reparametrization, (1.3) can be brought into the form (1.1) (but with W depending on more than one variable) for all gauge groups, and it was conjectured in [2] that this should describe the Seiberg-Witten effective action. String duality implies that this must indeed be true. However, for the groups $G = E_n, n \geq 6$, x_2 does not enter quadratically and thus cannot easily be integrated out (for example, one has $W_{E_6}^{\text{ALE}} = x_1^3 + x_2^4 + x_3^2 + \dots$). This means that for E_n , the manifolds (1.3) have *a priori* no obvious relation to Riemann surfaces, and appear intrinsically as higher dimensional surfaces.

The question then immediately arises as to how the manifolds (1.3) are related to the curves (1.2) for exceptional gauge groups. It is very natural to believe that somehow the periods must be the same, but how this precisely works was not clear until now. It is the purpose of this letter to show that from the point of view of SW theory, the curves (1.2) and the manifolds (1.3) are indeed physically equivalent.

There is also a physics aspect to this. In fact, (1.3) represents (a local, non-compact piece of) a threefold, and the BPS states correspond to wrappings of type IIB 3-branes [11] around the local vanishing homology H_3 . For $G = SU(n)$, it follows from the considerations in [9] that upon integrating out the quadratic terms in x_2 and x_3 , the wrapped 3-branes are *effectively* equivalent to wrapping 1-branes around the curve (1.2); the 1-branes, which are the left-over pieces of the 3-branes, are precisely the non-critical, (anti-)self dual strings of [12]. On physical grounds, one would expect this to be true also for the other, and in particular the exceptional, gauge groups, but this is at first sight not so obvious. By performing the integrals over x_2 and x_3 in $W_{E_6}^{\text{ALE}}(x_1, x_2, x_3; u_j)$, we will argue that indeed the 3-branes, when wrapped around the vanishing cycles of (1.3), are equivalent to 1-branes on the curves (1.2). We expect that a similar story holds for the other exceptional groups as well.

The potential importance of this type of considerations lies not in the generalization to other gauge groups of what has been known before, but ultimately in the generalization to other, non-field theory limits of type IIB strings. In a threefold, more general singularities than those describing YM theories typically do arise, and the question is under what conditions $\alpha' \rightarrow 0$ physics can effectively be described in terms of SW-like curves (with anti-self-dual strings wrapping around them).

2. Seiberg–Witten Curves

In [8] a general scheme was presented for obtaining a family of Seiberg–Witten Riemann surfaces for an $N = 2$ supersymmetric Yang–Mills theory with arbitrary gauge group, G . Specifically, given any representation \mathcal{R} of G , one considers the polynomial:

$$P_{\mathcal{R}}(x; u_j) = \det(x - \Phi_0) , \quad (2.1)$$

where Φ_0 is a generic adjoint Higgs v.e.v. written in the representation \mathcal{R} . One can diagonalize Φ_0 to some Cartan subalgebra element $v \cdot H$, where v is an ℓ -dimensional vector. In the representation \mathcal{R} , the eigenvalues of Φ_0 are $\lambda \cdot v$, where λ are the weights of \mathcal{R} . Thus one can write:

$$P_{\mathcal{R}} = \prod_{\lambda} (x - \lambda \cdot v) . \quad (2.2)$$

The polynomial, $P_{\mathcal{R}}(x; u_j)$, naturally decomposes into factors corresponding to Weyl orbits of weights. If the representation \mathcal{R} is miniscule, then by definition, the weights of \mathcal{R} form exactly one such orbit. If the representation is not miniscule, then we will take $P_{\mathcal{R}}$ to be any one of the Weyl orbit factors in the determinant (2.1). Let M be the degree of $P_{\mathcal{R}}$.

To get the Riemann surface for a simply laced group, G , one first defines a function $\tilde{P}_{\mathcal{R}}(x, \zeta; u_j, \mu)$ via

$$\tilde{P}_{\mathcal{R}}(x, \zeta; u_j, \mu) \equiv P_{\mathcal{R}}(x; u_1, \dots, u_{\ell-1}, u_{\ell} + \zeta + \mu^2/\zeta) . \quad (2.3)$$

That is, one shifts the top Casimir by $u_{\ell} \rightarrow u_{\ell} + \zeta + \mu^2/\zeta$. The Riemann surface is then given by

$$\tilde{P}_{\mathcal{R}}(x, \zeta; u_j, \mu) = 0 \quad (2.4)$$

The canonical way to view this surface [13,14,8] is as an M -sheeted foliation by $x(\zeta)$ over the base ζ -sphere. The sheets are then in one-to-one correspondence with the weights of \mathcal{R} (or at least the Weyl orbit of weights that one has chosen). There are $2(\ell + 1)$ branch points in the base, and they come in pairs related by the involution symmetry $\zeta \rightarrow \mu^2/\zeta$. The first ℓ pairs of branch points can be labelled by a system of simple roots $\{\alpha_j, j = 1, \dots, \ell\}$ of G , and the monodromy around the branch point is then given by the Weyl reflection r_{α_j} . Above the α_j -cut one then connects sheets in pairs according to whether the weights labelling the sheets are exchanged by the Weyl reflection r_{α_j} . The $(\ell + 1)$ -th pair of branch points is $\{\zeta = 0, \zeta = \infty\}$ and each of these points has the Coxeter monodromy corresponding to the product of the fundamental Weyl reflections. The sheets are joined above $\zeta = 0$ and $\zeta = \infty$ according to the Coxeter orbits of the weight labels.

The genus of the Riemann surface is usually far larger than the rank, ℓ , of G . Thus one needs to isolate a special sub-Jacobian, or, more precisely, a special Prym Variety, whose periods give the effective action of the $U(1)^\ell$ on the Coulomb branch. This can be done using the underlying integrable system, and can be implemented directly in a number of ways [13,14,8]. The simplest is to first take the ℓ cycles that surround the α_j -cuts on the ζ -sphere, and then take the inverse image under the projection of the Riemann surface to the ζ -sphere. This gives ℓ A -cycles. The B -cycles are then obtained by finding the cycles that intersect the A -cycles in the proper manner. As we will see later, once one has the Seiberg-Witten differential λ_{SW} , the issue of the proper A and B cycles is essentially moot.

Because of the connection with integrable systems, the Seiberg-Witten differential takes the universal form:

$$\lambda_{SW} = -2x \frac{d\zeta}{\zeta} . \quad (2.5)$$

We now focus on the details of the curve for E_6 . The simplest E_6 curve is obtained from the (miniscule) 27-dimensional representation. Given a system of simple roots, $\alpha_1, \dots, \alpha_6$, one finds that for each root, α_j , there are six² weights, $\lambda_j^{(a)}$, of the 27 such that $\lambda_j^{(a)} \cdot \alpha_j = +1$. Consequently, $\lambda_j'^{(a)} \equiv \lambda_j^{(a)} - \alpha_j$ are weights of the 27 with $\lambda_j'^{(a)} \cdot \alpha_j = -1$. Thus above *each* of the six α_j -cuts on the ζ -sphere, the sheets of the foliation are connected in six pairs, making a total of 36 such interconnections. Under the action of the Coxeter element the weights form three orbits of order 12, 12 and 3 respectively. If one imagines assembling the surface by first making the connections at $\zeta = 0$ and $\zeta = \infty$, one first gets three disjoint spheres (the Coxeter orbits) that must then be laced together by the 36 pieces of plumbing mentioned earlier. The result is a genus 34 surface.

This surface is explicitly given by

$$\frac{1}{2} x^3 \tau^2 - q_1 \tau + q_2 = 0 ; \quad \tau \equiv \zeta + \frac{\mu^2}{\zeta} + u_6 , \quad (2.6)$$

² The fact that this is the rank of E_6 is a coincidence.

where

$$\begin{aligned}
q_1 &= 270 x^{15} + 342 u_1 x^{13} + 162 u_1^2 x^{11} - 252 u_2 x^{10} + (26 u_1^3 + 18 u_3) x^9 \\
&\quad - 162 u_1 u_2 x^8 + (6 u_1 u_3 - 27 u_4) x^7 - (30 u_1^2 u_2 - 36 u_5) x^6 \\
&\quad + (27 u_2^2 - 9 u_1 u_4) x^5 - (3 u_2 u_3 - 6 u_1 u_5) x^4 - 3 u_1 u_2^2 x^3 \\
&\quad - 3 u_2 u_5 x - u_2^3 ; \\
q_2 &= \frac{1}{2x^3} (q_1^2 - p_1^2 p_2) ; \\
p_1 &= 78 x^{10} + 60 u_1 x^8 + 14 u_1^2 x^6 - 33 u_2 x^5 + 2 u_3 x^4 - 5 u_1 u_2 x^3 - u_4 x^2 \\
&\quad - u_5 x - u_2^2 ; \\
p_2 &= 12 x^{10} + 12 u_1 x^8 + 4 u_1^2 x^6 - 12 u_2 x^5 + u_3 x^4 - 4 u_1 u_2 x^3 - 2 u_4 x^2 \\
&\quad + 4 u_5 x + u_2^2 .
\end{aligned} \tag{2.7}$$

Note that (2.6) is *not* a hyperelliptic curve, unlike the simplest curves for the A_n and D_n groups. Moreover, like the curves for D_n , the genus of the curve exceeds the rank of the gauge group, but unlike D_n , there is no obvious, elementary symmetry that picks out the cycles that yield the quantum effective action [5]. For the curve defined by (2.6) one has to use the methods outlined above and in [8] to determine the cycles of interest.

It is useful to note that since (2.6) is quadratic in τ , one can solve it to obtain a more convenient presentation of the curve:

$$\tau = \zeta + \mu^2/\zeta + u_6 = \frac{1}{x^3} [q_1 \pm p_1 \sqrt{p_2}] , \tag{2.8}$$

where q_1, p_1 and p_2 are defined above.

This expression for τ will be of importance in the next section, but it is also of interest for other reasons. One may view (2.8) as determining how each eigenvalue, x , of the v.e.v., Φ_0 , changes as one varies the top Casimir, u_6 , while holding the remaining Casimirs, u_1, \dots, u_5 fixed. As a result, the explicit expression (2.8), which has degree 12, can be interpreted as a single variable (though non-polynomial) version of the Landau-Ginzburg potential for E_6 , in the sense of [15]. It can probably be used to describe the coupling of E_6 minimal matter to topological gravity directly, using the residue methods of [16,15]³.

³ Indeed, there was a long-standing question how the KdV type Lax operator, which is derived from the 27-dimensional representation of E_6 via Drinfeld-Sokolov reduction and which thus has *a priori* degree 27, is related to the E_6 simple singularity of degree 12. This simple singularity supposedly figures as the superpotential of a topological LG theory that describes the matter-gravity system [16].

Finally, we observe that one can also obtain (2.8) from decomposing Casimir invariants of E_6 into those of $SO(10) \times U(1)$. If one writes the u_j in terms of the five invariants of $SO(10)$ and the $U(1)$ eigenvalue x , and then inverts this relationship, one can easily extract an expression for u_6 in terms of u_1, \dots, u_5 and x . Thus one can easily derive (2.8) from the results of [17]. That this is equivalent to the procedure described above follows from the fact that the stabilizer of a weight space in the 27 of E_6 is $SO(10) \times U(1)$.

3. The quantum effective action from string theory

It was argued in [10,18,9] that one can obtain the quantum effective action of a pure gauge theory from the IIB string compactified on a Calabi-Yau manifold that is degenerating to a nearly singular ALE fibration over a \mathbb{P}^1 base. The Riemann surface can then be constructed⁴ by using the monodromy data on the non trivial 2-cycles of the fiber to produce the monodromies of the Riemann sheets. The Seiberg-Witten differential is obtained by integrating the holomorphic $(3,0)$ -form, Ω , of the Calabi-Yau manifold over the 2-cycles in the fiber, yielding a meromorphic 1-form on the Riemann surface [9]. This is easily verified for A_n gauge groups,⁵ but it is far less obvious for the exceptional groups. Here we will describe, in detail, how it works for E_6 .

The singular ALE fibration for E_6 has the form:

$$\mathcal{P}(y_1, y_2, y_3; \zeta) \equiv W_{E_6}(y_1, y_2, y_3) + \nu (\zeta + \mu^2/\zeta) = 0, \quad (3.1)$$

where ζ is the coordinate on the base \mathbb{P}^1 and

$$\begin{aligned} W_{E_6}(y_1, y_2, y_3) = & y_1^3 + y_2^4 + y_3^2 + w_1 y_1 y_2^2 + w_2 y_1 y_2 + w_3 y_2^2 \\ & + w_4 y_1 + w_5 y_2 + w_6. \end{aligned} \quad (3.2)$$

The parameter ν in (3.1) is a normalization constant that has been inserted for later convenience. The holomorphic $(3,0)$ -form in these local coordinates can be written:

$$\Omega = \left(\frac{d\zeta}{\zeta} \right) \wedge \frac{dy_1 \wedge dy_2}{y_3} = \left(\frac{d\zeta}{\zeta} \right) \wedge \frac{dy_1 \wedge dy_2}{\sqrt{\mathcal{P}(y_1, y_2, y_3 = 0; \zeta)}}. \quad (3.3)$$

⁴ Note that, strictly speaking, the SW curve itself is not geometrically embedded in the Calabi-Yau manifold.

⁵ It is straightforward to extend the arguments of [9] to D_n gauge groups as well.

The appearance of $d\zeta/\zeta$ in (3.3) follows from the Ricci flatness of the underlying Calabi-Yau and the Ricci flatness of the fiber.

We need to find the 2-cycles in the fiber space defined by $W_{E_6} = \text{const}$, and then we need to integrate (3.3) over these fibers. The key to doing this is to recall some beautiful facts of classical algebraic geometry [19]. One starts by recasting (3.2) as a cubic in \mathbb{P}^3 [20]. That is, given

$$\begin{aligned} W(x_i) = & x_3^2 x_4 + x_1^3 + 2i x_2^2 x_3 + w_1 x_1 x_2^2 + w_2 x_1 x_2 x_4 + w_3 x_2^2 x_3 \\ & + w_4 x_1 x_4^2 + w_5 x_2 x_4^2 + w_6 x_4^3, \end{aligned} \quad (3.4)$$

in homogeneous coordinates, x_i , one can obtain (3.2) by going to the patch $x_4 \neq 0$, setting $y_1 = x_1/x_4$, $y_2 = x_2/x_4$ and $y_3 = (x_3/x_4 + ix_2^2/x_4^2)$. A cubic in \mathbb{P}^3 can be thought of as the blow-up of six points in general position in \mathbb{P}^2 [19]. The second homology, H_2 , of this surface is seven dimensional, and consists of these six (non-intersecting) spheres, and the canonical class of the \mathbb{P}^2 , which has intersection number 1 with each of the six blown up spheres. The six dimensional integer homology of the local E_6 singularity consists on a six-dimensional subspace of H_2 . This subspace is obtained by orthogonalizing with respect to a vector with coordinates $(3, -1, -1, -1, -1, -1)$, and the result is a set of homology cycles that have the intersection matrix of E_6 [19,20,13].

Thus far, we have seen very little that looks like the 27-sheeted Riemann surface described in the previous section. However, a celebrated fact about a cubic in \mathbb{P}^3 is that it contains 27 lines. That is, there are 27 holomorphically embedded \mathbb{P}^1 's in this space. Each of these 27 lines is a non-trivial element of homology, and together they (over) span the (seven dimensional) homology. Thus, we can exhibit the homology by exhibiting these lines. Moreover, it is also well known that the monodromy group of the E_6 singularity is the Weyl group of E_6 , and that it acts on these lines as on the 27 of E_6 ⁶.

The explicit computation of the 27 lines for the generic E_6 singularity was recently given by Minahan and Nemeschansky [21], and while the motivation of these authors was rather different from ours, our analysis will closely parallel theirs. Make a change of variables from (y_1, y_2) to (x, y) in (3.2), where $y_1 = x y + \alpha(x)$, and $y_2 = y + \beta(x)$. As yet, α and β are arbitrary functions of x . With these changes of variable, the function \mathcal{P}

⁶ As a representation of the Weyl group of E_6 , the 27 is reducible: $27 = 20 + 6 + 1$. The singlet is the cycle with respect to which one orthogonalizes to get the cycles of the E_6 singularity, and the 6 is the fundamental reflection representation (the Cartan subalgebra representation) on the compact homology basis of the E_6 singularity.

in (3.1) is a quartic in y . Choosing $\beta = -(x^3 + w_1 x)/4$ cancels the y^3 term in this quartic. One next chooses $\alpha(x)$ so as to get rid of the linear term in y . This involves solving a quadratic equation for $\alpha(x)$, however, before doing this, we wish to reparametrize the versal deformation of the E_6 singularity via:

$$\begin{aligned} w_1 &= \frac{1}{2} u_1 ; & w_2 &= -\frac{1}{4} u_2 ; & w_3 &= \frac{1}{96} (u_3 - u_1^3) ; \\ w_4 &= \frac{1}{96} (u_4 + \frac{1}{4} u_1 u_3 - \frac{1}{8} u_1^4) ; & w_5 &= -\frac{1}{48} (u_5 - \frac{1}{4} u_1^2 u_2) ; \\ w_6 &= \frac{1}{3456} (u_6 + \frac{1}{16} u_1^6 - \frac{3}{16} u_1^3 u_3 + \frac{3}{32} u_3^2 - \frac{3}{4} u_1^2 u_4) . \end{aligned} \quad (3.5)$$

One then finds that

$$\alpha(x) = \frac{1}{48x} \left[(2u_1 x^3 + u_1^2 x + 2u_2) \pm 2 \sqrt{p_2} \right] . \quad (3.6)$$

where p_2 is given by (2.7).

The lines in cubic occur when one can analytically solve $\mathcal{P} = 0$. Since \mathcal{P} is quadratic in y_3 , this happens precisely when the rest of \mathcal{P} is a perfect square. That is, we must find the points at which $\mathcal{P}(y_1, y_2, y_3 = 0; \zeta)$ becomes a perfect square. With the changes of variable above, and the choices of $\alpha(x)$ and $\beta(x)$, the function $\mathcal{P}(x, y, y_3 = 0)$ takes the form of a quadratic in y^2 . The lines are thus defined by the vanishing of the discriminant, Δ , of this quadratic in y^2 . A straightforward computation shows that:

$$\Delta = \frac{1}{864} \left[\frac{q_1 \mp p_1 \sqrt{p_2}}{x^3} - u_6 \right] - 4 \nu (\zeta + \mu^2/\zeta) . \quad (3.7)$$

Taking $\nu = \frac{1}{3456}$, one thus finds that is discriminant vanishes precisely when x and ζ satisfy the equation (2.6). We thus see that the Riemann surface is nothing other than a fibration, over the ζ -sphere, of the locations of the 27 lines in the E_6 singularity.

The lines are all defined by a relationship of the form:

$$y_3^2 = -(y^2 + a(x))^2 , \quad (3.8)$$

and hence, from the perspective of the singularity polynomial, these lines are all non-compact cycles in the closed homology of singularity relative to the boundary of the singularity [6]. Note that the coefficient of y^2 in (3.8) is independent of x , and so all the lines meet at infinity. This means that the difference of any two lines defines a compact cycle of the singularity. It is over these compact cycles that we want to integrate Ω .

With the change of variables above, Ω now takes the form

$$\Omega = \left(\frac{d\zeta}{\zeta} \right) \wedge \frac{(y + \beta'(x) - x\alpha'(x)) dx \wedge dy}{\sqrt{y^4 + b_1(x, \zeta; u_j) y^2 + b_0(x, \zeta; u_j)}} \quad (3.9)$$

At generic values of x and ζ there are two branch cuts for Ω (related by $y \rightarrow -y$) in the y -plane. These cuts simultaneously disappear at each of the 27 lines. As in [9], the “latitude circles” of the homology 2-cycles correspond to contours that surround the branch cuts, and these “latitude circles” vanish at the North and South poles of the 2-cycles defined by the vanishing of the branch cuts at the lines. In this way, a pair of lines defines a homology 2-cycle. The only difference here is that there are now two branch cuts (as opposed to one such cut for A_n and D_n), and this introduces a minor subtlety.

Given two lines at which the cuts simultaneously vanish, there are two ways in which the branch points can behave. Suppose that near one line, L , the branch points appear in pairs $\{y = \xi, y = \eta\}$, and $\{y = -\xi, y = -\eta\}$ with $\xi \rightarrow \eta$ as one approaches L . Take the cuts, C_{\pm} , to run from $\pm\xi$ to $\pm\eta$ respectively. As one approaches another line, L' , either $\xi \rightarrow \eta$ again, or $\xi \rightarrow -\eta$. For the former possibility, a contour surrounding either cut defines a non-trivial 2-cycle, and the two 2-cycles defined by the two cuts do not intersect. If, however, one has $\xi \rightarrow -\eta$, then the cuts C_+ and C_- annihilate one another. Each cut then defines the half of a 2-cycle, and the annihilation of the cuts corresponds to gluing hemispheres of the 2-cycle together. These two possibilities are reflected in the fact that the difference of two different weights in the 27 of E_6 can either have length-squared 4 or 2. For the former, the difference weights is the sum of two orthogonal roots of E_6 , and the latter corresponds to a single root of E_6 . It follows that if the lines L and L' give rise to a pair of non-intersecting 2-cycles, then there must be another intermediate line, L'' , such that paths from L to L'' and from L'' to L' describe each of the spheres separately.

It should now be evident that the proper contour over which to integrate Ω is one that surrounds both cuts, and that does not get pinched off as the cuts annihilate. This simple closed loop in the y -plane can then be deformed to large $|y|$, and the integral $\oint \Omega dy$ becomes the elementary integral: $\oint \frac{dy}{y} = 2\pi i$. The integral over x is now trivial: $2\pi i \int dx$ evaluated between the lines defining the poles and equatorial gluing of hemispheres of the 2-cycle in question. One thus finds that the integral of Ω over the 2-cycle, L_{ij} , defined by the lines L_i and L_j is:

$$\int_{L_{ij}} \Omega = 2\pi i (x_j - x_i) \left(\frac{d\zeta}{\zeta} \right) = -\pi i (\lambda_{SW}|_{x=x_j} - \lambda_{SW}|_{x=x_i}), \quad (3.10)$$

where x_i and x_j are the locations of the lines. Thus we recover the difference of the Seiberg-Witten differential between the sheets of the Riemann surface, exactly as in [9].

Given that there are really only six compact 2-cycles in the fiber of the foliation, it follows that there are 12 compact 3-cycles (six A -cycles and six B -cycles) in the total

space of the foliation (3.1). Recalling the discussion of the branch cuts in the base, and the plumbing of the Riemann surface, we see that above each α_j -cut, there are six copies of the α_j 2-cycle expanding from a branch point and collapsing back to a branch point. The result is the A -type 3-cycle α_j . (The B -type cycles are made by going from one of these branch points out to $\zeta = 0$ or $\zeta = \infty$.) Thus the redundancy of the genus 34 Riemann surface stems from the highly redundant description of a six-dimensional fiber homology in terms of 27 lines.

This establishes a useful result about the integrals of λ_{SW} on the Riemann surface (2.6): since there are only six A -type 3-cycles and six B -type 3-cycles, there can only be six independent integrals of Ω over the sets of A and B cycles. This means that there can only be six independent integrals of λ_{SW} over the A -cycles of the Riemann surface, and six independent integrals over B -cycles, even though the genus of the surface is very large. Moreover, a basis of such integrals can be obtained by choosing *any* representative cycle in the surface that lies above each of the α_j cuts.

4. Comments

The lines on the surface in projective space, or the non-compact cycles of the local singularity type, have played a central role in defining the sheets of the Riemann surface. This was also true of the analysis in [9], though it was not explicitly stated there. The fact that the lines in the projective space span a vector space of dimension $\ell + 1$ suggests that there may also be a further 2-cycle or 3-cycle in the Calabi-Yau manifold that is playing an interesting implicit role.

While we have treated the E_6 theory in detail here, there is very probably a similar story for the other E_ℓ groups. The corresponding singularities can all be realized as the blow-up of ℓ points in general position ($\ell \leq 8$) in a \mathbb{P}^2 [13]. One can then presumably realize these surfaces in some weighted projective spaces, and then use the lines on these surfaces to relate the integration of Ω over 3-cycles to the integration of λ_{SW} over cycles of Riemann surfaces, in analogy to what we did above.

In [21] the E_6 singularity and the 27 lines were used to construct a candidate Seiberg-Witten curve for a superconformal theory with matter and E_6 global symmetry. In string theory, global symmetries are generically gauged [22], and so one might expect some connection between the results here and those of [21] via some kind of flavour gauging. It turns out that we can come fairly close: there are two extreme degenerations of (2.6) at which the theory becomes superconformal: (i) the E_6 Argyres-Douglas points [23,24], and (ii)

the point where the Argyres-Douglas points come together: $\mu = 0$, *i.e.* when $\tau = \zeta + u_6$. In the latter instance, the Riemann surface collapses to a genus 4 surface foliated over the x -sphere:

$$\left[\frac{x^3 (\zeta + u_6) - q_1}{p_1} \right]^2 = p_2 . \quad (4.1)$$

The Seiberg-Witten differential, (2.5), is holomorphic on this surface except at infinity and at the zeroes of ζ . At infinity the differential has a double pole, while at the zeroes of ζ it has a simple pole with residue $-2x$. From (2.2), (2.3) and (2.4) one immediately sees that there are 27 such points on the base x -sphere, and the residues are simply $-2\lambda \cdot v$. If one views the six-dimensional vector v as the masses of some matter fields, then one has obtained a theory with matter with an E_6 global symmetry. The most significant physical difference between this matter theory, and that of [21], is that this “matter theory” lacks an additional free parameter, called ρ in [21]. This parameter describes a Higgs v.e.v. on the Coulomb branch, but in our degeneration limit, the Coulomb branch has been collapsed to provide the masses.

In summary, we have made precise (for E_6) something that was conjectured some while ago [2]: namely that the *ADE* singularity types must be the key ingredients in the construction of the quantum effective actions of the $N = 2$ supersymmetric Yang-Mills theories with *ADE* gauge groups. As in [9] we have also established an important string theory result: the E_6 Yang-Mills theory can be represented in terms of compactifying the six-dimensional self-dual string [12] on the genus 34 Riemann surface defined by (2.6).

Acknowledgements

We would like to thank the Aspen Center for Physics for providing a stimulating research atmosphere and the opportunity to complete this work. We would like to thank P. Candelas, T. Eguchi, B. Greene, K. Hori, D. Nemeschansky, and particularly E. Martinec, S. Katz and D. Morrison for valuable discussions. Some of the details of the Riemann surface described in section 2 were done as part of a collaboration with E. Martinec [8], but were not published. N.W. is also grateful to the theory division at CERN for hospitality while some of this work was undertaken. N.W. is supported in part by funds provided by the DOE under grant number DE-FG03-84ER-40168.

References

- [1] N. Seiberg and E. Witten, Nucl. Phys. B426 (1994) 19, hep-th/9407087; Nucl. Phys. B431 (1994) 484, hep-th/9408099.
- [2] A. Klemm, W. Lerche, S. Theisen and S. Yankielowicz, Phys. Lett. B344 (1995) 169, hep-th/9411048.
- [3] P. Argyres and A. Faraggi, Phys. Rev. Lett. **74** (1995) 3931, hep-th/9411057.
- [4] U. Danielsson and B. Sundborg, Phys. Lett. BB358 (1995) 273, hep-th/9504102.
- [5] A. Brandhuber and K. Landsteiner, Phys. Lett. BB358 (1995) 73, hep-th/9507008.
- [6] V. Arnold, A. Gusein-Zade and A. Varchenko, *Singularities of Differentiable Maps II*, Birkhäuser 1985.
- [7] A. Gorskii, I. Krichever, A. Marshakov, A. Mironov and A. Morozov, Phys. Lett. B355 (1995) 466, hep-th/9505035;
H. Itoyama and, A. Morozov, preprints ITEP-M5-95 and ITEP-M6-95, hep-th/9511126 and hep-th/9512161.
- [8] E. Martinec and N.P. Warner, Nucl. Phys. B459 (1996) 97, hep-th/9509161
- [9] A. Klemm, W. Lerche, P. Mayr, C. Vafa and N.P. Warner, hep-th/9604034.
- [10] S. Kachru and C. Vafa, Nucl. Phys. B450 (1995) 69, hep-th/9505105.
- [11] A. Strominger, Nucl. Phys. B451 (1995) 96, hep-th/9504090.
- [12] M. Duff and J. Lu, Nucl. Phys. B416 (1994) 301, hep-th/9306052;
E. Witten, preprint IASSNS-HEP-95-63, hep-th/9507121;
A. Strominger, hep-th/9512059;
O. Ganor and A. Hanany, preprint IASSNS-HEP-96-12, hep-th/9602120;
N. Seiberg and E. Witten, preprint RU-96-12, hep-th/9603003;
M. Duff, H. Lu and C.N. Pope, preprint CTP-TAMU-9-96, hep-th/9603037.
- [13] V. Kanev, *Spectral curves, simple Lie algebras, and Prym-Tjurin varieties*, in *Theta Functions, Bowdoin 1987*, Proc. Symp. Pure Math. 49, L. Ehrenpreis and R. Gunning (eds.), AMS, Providence (1989).
- [14] R. Donagi, *Decomposition of spectral covers*, in *Journées de Geometrie Algébrique d'Orsay*, Asterisque vol. 218, Soc. Math. de France, Paris (1993).
- [15] B. Dubrovin, *Geometry of 2d Topological Field Theories*, SISSA-89/94/FM;
Integrable Systems in Topological Field Theory, Nucl. Phys. B379 (1992) 627-689;
Hamiltonian Formalism of Witham-type Hierarchies in Topological Landau-Ginzburg Model, Comm. Math. Phys. **145** (1992) 19.

- [16] R. Dijkgraaf, E. Verlinde and H. Verlinde, Nucl. Phys. B352 (1991) 59.
- [17] W. Lerche and N.P. Warner, Nucl. Phys. B358 (1991) 571.
- [18] S. Kachru, A. Klemm, W. Lerche, P. Mayr and C. Vafa, Nucl. Phys. B459 (1996) 537, hep-th/9508155.
- [19] R.Hartshorne, *Algebraic Geometry*, Springer-Verlag (1977).
- [20] S. Katz and D. Morrison, Jour. Alg. Geom. 1 (1992) 449.
- [21] J. Minahan and D. Nemeschansky, hep-th/9608047.
- [22] T. Banks, M. Douglas and N. Seiberg, preprint RU-96-41, hep-th/9605199.
- [23] P. Argyres and M. Douglas, Nucl. Phys. B448 (1995) 93, hep-th/9505062.
- [24] T. Eguchi and K. Hori, preprint UT-755, hep-th/9607125.